

Probing the Origin of the EMC Effect via Tagged Structure Functions of the Deuteron

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Abstract

We demonstrate that measurement of tagged structure functions of the deuteron in $(e, e'N)$ semi-inclusive reactions can discriminate between different hypotheses on the origin of the nuclear EMC effect. By choosing extreme backward kinematics for the spectator nucleon to minimize effects from the deuteron wave function and final state interactions, one can isolate the modifications in the structure of the bound nucleon within the impulse approximation. The same reaction can be used to extract the large- x neutron to proton structure function ratio.

I. INTRODUCTION

More than a decade after the discovery of the nuclear EMC effect [1] and many fine measurements [2–5] of the ratios of structure functions of nuclei and the deuteron, no consensus has been reached on the origin of the effect. The x dependence of the effect, while non-trivial, is rather smooth and has the same basic shape for all nuclei, making it is easy to fit in a wide range of models with very different underlying assumptions. The only extra constraint available so far comes for measurements of the A -dependence of the sea distribution, which restricts some of the models, but is still not sufficient to allow one to unambiguously identify the origin of the EMC effect.

In order to move beyond this rather unsatisfactory situation, new experiments involving more kinematical variables accessible to accurate measurements are necessary. The aim of this study is to demonstrate that use of semi-inclusive processes off the deuteron,

$$e + D \rightarrow e + N + X, \quad (1)$$

where a nucleon is detected in the target deuteron fragmentation region, may help to discriminate between some classes of models. In particular, one may be able to distinguish between models in which the effect arises entirely from hadronic degrees of freedom — nucleons and pions (which in the traditional nuclear physics picture are responsible for the binding of the nucleus), and models in which the effect is attributed to the explicit deformation of the wave function of the bound nucleon itself. By selecting recoil nucleons with small transverse momentum in the backward region, effects due to final state interactions (FSI) can be minimized, thus allowing one to probe, within the impulse approximation, the deformation of the structure of the bound nucleon.

It is worth emphasizing that measurement of the reaction (1) in the kinematics of interest will become feasible in the near future, for example at Jefferson Lab [6]. Recent development of silicon detectors also make it possible in the jet target experiments to measure recoil nucleons with a low momentum threshold of about 100–150 MeV/c. Plans to implement such a technique are under discussion at the HERMES detector at HERA [7].

Aside from providing insight into the origin of the nuclear EMC effect per se, the measurements of tagged events may also be useful in connection with the problem of extracting the neutron structure function from deuteron data. By selecting only the slowest recoil protons in the target fragmentation region, one should be able to isolate the situation whereby the virtual photon scatters from a nearly on-shell neutron in the deuteron. In this way one may hope to extract the F_{2n} structure function while minimizing uncertainties arising from modeling the nuclear effects in the deuteron.

This paper is organized as follows. In Section II we outline the basic formalism for semi-inclusive deep-inelastic scattering off the deuteron, and discuss the conditions under which the impulse approximation may be valid. Section III is devoted to a survey of several models whose predictions for the ratios of tagged structure functions are compared. Possibilities of extracting the neutron structure function from semi-inclusive experiments are discussed in Section IV, and finally some conclusions are drawn in Section V.

II. BASIC FORMALISM

It was observed a long time ago that the nuclear EMC effect for the deviation from unity of the ratio

$$R(x, Q^2) = \frac{2F_{2A}(x, Q^2)}{AF_{2D}(x, Q^2)} \quad (2)$$

is approximately proportional to the nuclear density. This is natural for a dilute system, and indicates that most of the EMC effect is due to two-nucleon interactions. Based on this observation one should expect that the EMC effect in the deuteron is much smaller than that in heavy nuclei. However, by virtue of the uncertainty principle, one may try to enhance the effect by isolating the configuration where the two nucleons in the deuteron are close together. For example, it is easy to check that the main contributions to the deuteron wave function for nucleon momenta $\mathbf{p} \gtrsim 300$ MeV/c come from distances $\lesssim 1.2$ fm.

To make our discussion quantitative, let us begin by writing the electromagnetic tensor $W_D^{\mu\nu}$ of the deuteron in terms of the matrix element of the electromagnetic current:

$$W_D^{\mu\nu} = \sum_{spin, X} \langle D | J_D^\mu(q) | X, N \rangle \langle X, N | J_D^{\mu\dagger}(0) | D \rangle. \quad (3)$$

The tensor $W_D^{\mu\nu}$ can be expanded in terms of four possible Lorentz structures, with the coefficients given by the invariant structure functions F_L^D, F_T^D, F_{TL}^D and F_{TT}^D (T = transverse, L = longitudinal). These functions can in general depend on four variables, constructed from the four-momenta of the target deuteron (P), virtual photon (q) and spectator nucleon (p_s) (or equivalently the momentum of the struck nucleon p , where $\mathbf{p} = -\mathbf{p}_s$ in the deuteron rest frame). In terms of the invariant structure functions the differential cross section for the semi-inclusive reaction ($e, e'N$) can then be written:

$$\begin{aligned} \frac{d\sigma}{dx dQ^2 d^3\mathbf{p}_s/E_s} &= \frac{2\alpha_{em}^2}{xQ^4} \left(1 - y - \frac{x^2 y^2 M^2}{Q^2} \right) \\ &\times \left[F_L^D + \left(\frac{Q^2}{2\mathbf{q}^2} + \tan^2(\theta/2) \right) \frac{\nu}{M} F_T^D + \left(\frac{Q^2}{\mathbf{q}^2} + \tan^2(\theta/2) \right)^{\frac{1}{2}} \cos\phi F_{TL}^D + \cos(2\phi) F_{TT}^D \right], \quad (4) \end{aligned}$$

where the structure functions are related to the components of the electromagnetic tensor $W_D^{\mu\nu}$ by:

$$F_L^D = \nu(1 + \cos\delta)^2 \cdot W_D^{--}, \quad (5a)$$

$$F_T^D = M(W_D^{xx} + W_D^{yy}), \quad (5b)$$

$$F_{TL}^D = 2\nu(1 + \cos\delta) \cdot W_D^{-x}, \quad (5c)$$

$$F_{TT}^D = \frac{\nu}{2} \sin^2\delta \cdot (W_D^{xx} - W_D^{yy}). \quad (5d)$$

The kinematic variables in Eqs.(4) and (5) are $x = Q^2/2M\nu$, where ν is the energy of virtual photon in the target rest frame and M the nucleon mass, $Q^2 = 4E(E - \nu) \sin^2(\theta/2)$ is the squared four-momentum transfer to the target, E is the beam energy and $y = \nu/E$. The angle ϕ is the azimuthal angle for the spectator nucleon, and $E_s = \sqrt{M^2 + \mathbf{p}_s^2}$ the spectator nucleon energy, and $\sin^2\delta = Q^2/\mathbf{q}^2$ in Eqs.(5). We define the photon three-momentum \mathbf{q} to be in the $+z$ direction.

A convenient choice of the four independent variables for the structure functions is the two inclusive deep-inelastic scattering variables, x and Q^2 , and the transverse momentum, \mathbf{p}_T^s , and light-cone momentum fraction, α_s , of the spectator nucleon:

$$\alpha_s = \frac{E_s - p_z^s}{M}, \quad (6)$$

where p_z^s is the longitudinal momentum of the detected nucleon (to simplify the expressions we have neglected here the deuteron binding energy, $M_D \approx 2M$). In the Bjorken limit the variable α_s then satisfies the condition [8]:

$$\alpha_s \leq 2 - x. \quad (7)$$

Having defined the relevant cross sections, we must now establish the kinematical range in which the nuclear modifications of the bound nucleon structure function will be accessible experimentally. From Eqs.(4) and (5) the restrictions:

$$Q^2/\mathbf{q}^2 \sim 4M^2x/Q^2 \ll 1, \quad (8a)$$

$$Q^2 \ll E(E - \nu), \quad (8b)$$

enhance the contribution of the longitudinal structure function F_L , which is expressed through the “good” component of the electromagnetic current, Eq.(5a). Another simplification can be achieved by considering the situation where the detected nucleon in the deuteron is in the spectator kinematics, namely:

$$\alpha_s \geq 1 - x. \quad (9)$$

In this kinematical region the contribution of the direct process where a nucleon is produced at the γ^*N interaction vertex is negligible [8].

A. Factorization and the Impulse Approximation

In formulating deep-inelastic scattering from the deuteron the simplest approach adopted has been the impulse approximation for the nuclear system. We will consider two formulations of the impulse approximation: one based on the covariant Feynman, or instant-form, approach, where one nucleon is on-mass-shell and one off-mass-shell, and the non-covariant/Hamiltonian light-cone approach, in which both nucleons are on-mass-shell (but off the light-cone energy shell). In the next subsection we shall consider corrections to the impulse approximation, in the form of final state interactions, but for now let us review briefly the basic impulse approximation assumptions and results.

1. Covariant Instant-Form Approach

In order to construct covariant amplitudes in the instant-form of quantization, summation over all possible time-orderings of intermediate states is essential. Incorporation of negative energy configurations into the total Lorentz-invariant amplitude is done on the basis of introducing intermediate state particles which are off their mass shells. In the impulse approximation for the scattering of a virtual photon, γ^* , from the deuteron, the invariant amplitude for the complete process factorizes into a product of amplitudes for γ^* -off-mass-shell nucleon scattering, and for forward nucleon-deuteron scattering. In this case the hadronic tensor of the deuteron can be written [9]:

$$W_D^{\mu\nu}(P, p, q) = \text{Tr} \left[\hat{\mathcal{S}}_{ND}(P, p) \hat{W}_N^{\mu\nu}(p, q) \right], \quad (10)$$

where the nucleon-deuteron scattering amplitude $\hat{\mathcal{S}}_{ND}$ (for spin-averaged processes) contains scalar and vector components:

$$\hat{\mathcal{S}}_{ND} = \hat{\mathcal{S}}_0 + \gamma_\alpha \hat{\mathcal{S}}_1^\alpha. \quad (11)$$

The operator $\hat{W}_N^{\mu\nu}$ in Eq.(10) is the truncated hadronic tensor for the off-shell nucleon, which describes the γ^*N interaction. Because $\hat{W}_N^{\mu\nu}$ is a matrix in both Lorentz and Dirac spaces, its structure is necessarily more general than that for a free nucleon. In the Bjorken limit, it was shown in Ref. [9] that out of a possible 13, only three independent structures contribute:

$$\widehat{W}_N^{\mu\nu}(p, q) = - \left(g_{\mu\nu} + \frac{q_\mu q_\nu}{Q^2} \right) \left(\widehat{W}_0(p, q) + \not{p} \widehat{W}_1(p, q) + \not{q} \widehat{W}_2(p, q) \right) + \mathcal{O}\left(\frac{1}{Q^2}\right). \quad (12)$$

Note that the full expression for $\widehat{W}_N^{\mu\nu}$ contains both positive and negative energy pieces. The on-shell nucleon tensor is obtained from $\widehat{W}_N^{\mu\nu}$ by projecting out the positive energy components:

$$W_N^{\mu\nu}(p, q) = \frac{1}{4} \text{Tr} \left[(\not{p} + M) \widehat{W}_{\mu\nu}(p, q) \right] \quad (13a)$$

$$= - \left(g_{\mu\nu} + \frac{q_\mu q_\nu}{Q^2} \right) \left(M \widehat{W}_0 + M^2 \widehat{W}_1 + p \cdot q \widehat{W}_2 \right), \quad (13b)$$

where the functions $\widehat{W}_{0\dots 2}$ here are evaluated at their on-shell points.

In terms of the off-shell truncated functions $\widehat{W}_{0\dots 2}$ the deuteron tensor can then be written:

$$W_D^{\mu\nu}(P, p, q) = -4 \left(g_{\mu\nu} + \frac{q_\mu q_\nu}{Q^2} \right) \left(\widehat{\mathcal{S}}_0 \widehat{W}_0 + \widehat{\mathcal{S}} \cdot p \widehat{W}_1 + \widehat{\mathcal{S}} \cdot q \widehat{W}_2 \right), \quad (14)$$

so that in general the right-hand-side does not factorize into a single term with separate nuclear and nucleon components. As explained in Ref. [9], factorization can be recovered by projecting from $\widehat{\mathcal{S}}_{0,1}$ only the positive energy (on-shell) components, in which case:

$$\widehat{\mathcal{S}}_0 \propto \widehat{\mathcal{S}} \cdot p \propto \widehat{\mathcal{S}} \cdot q. \quad (15)$$

This proportionality can also be obtained by taking the non-relativistic limit for the ND amplitudes to order \mathbf{p}^2/M^2 [10], since all negative energy contributions enter at higher orders [11,12]. For the relativistic ND interaction, the negative energy components enter through the P -state wave functions of the deuteron. However, in most realistic calculations of these, in the context of relativistic meson-exchange models of the deuteron [13], the contribution of the P -state wave functions is only a fraction of a percent, and so long as one avoids regions of extreme kinematics which are sensitive to the very large momentum components of the deuteron wave function, dropping the negative energy components is a reasonable approximation. Indeed, the explicit calculations of the inclusive structure function of the deuteron with relativistic wave functions shows that the factorization-breaking corrections amount to $\lesssim 1\text{--}2\%$ for all values of $x \lesssim 0.9$ [11,14].

Finally, given these approximations one can write the deuteron hadronic tensor as:

$$W_D^{\mu\nu}(P, p, q) \approx S^{IF}(P, p) W_N^{\mu\nu \text{ eff}}(p, q), \quad (16)$$

where S^{IF} is the nucleon spectral function within the instant-form impulse approximation [11]:

$$S^{IF}(\alpha_s, p_T) = (2 - \alpha_s) \frac{E_s M_D}{2(M_D - E_s)} |\psi_D(\alpha_s, p_T)|^2, \quad (17)$$

normalized such that

$$\int d^2\mathbf{p}_T \frac{d\alpha_s}{\alpha_s} S^{IF}(\alpha_s, p_T) = 1 \quad (18)$$

to ensure baryon number conservation [15,16]. Note that the approximation for the spectral function in Eq.(17) is valid only for non-relativistic momenta, and for large momenta the full (non-convolution) expression for the hadronic tensor in Ref. [11] should be used. The effective nucleon hadronic tensor $W_N^{\mu\nu \text{ eff}}$ is defined as:

$$\begin{aligned} W_N^{\mu\nu \text{ eff}} &= - \left(g_{\mu\nu} + \frac{q_\mu q_\nu}{Q^2} \right) \left(M \widehat{W}_0 + p^2 \widehat{W}_1 + p \cdot q \widehat{W}_2 \right) \\ &\equiv - \left(g_{\mu\nu} + \frac{q_\mu q_\nu}{Q^2} \right) F_{1N}^{\text{eff}} \left(\frac{x}{2 - \alpha_s}, p^2, Q^2 \right) + \dots, \end{aligned} \quad (19)$$

where the effective nucleon structure function F_{1N}^{eff} is now a function of the momentum fraction $x/(2 - \alpha_s)$ and the virtuality p^2 of the bound nucleon, as well as Q^2 . The kinematics of the spectator process gives rise to the relation:

$$p^2 = - \frac{2p_T^2 + (2 - \alpha_s)M^2}{\alpha_s} + \frac{1}{2}(2 - \alpha_s)M_D^2, \quad (20)$$

using which one can equivalently express F_{1N}^{eff} as a function of the transverse momentum p_T of the interacting nucleon, rather than p^2 .

2. Light-Cone Approach

In the light-cone formalism one can formally avoid the problems associated with negative energy solutions in Eq.(12), however, to obtain Eq.(16) one must on the other hand consider contributions arising from instantaneous interactions [17]. To take these effects into account one has to use gauge invariance to express the contribution of the “bad” current components of the electromagnetic tensor through the “good” components:

$$J_+^A = - \frac{q_+}{q_-} J_-^A, \quad (21)$$

and include the contribution of the instantaneous exchanges using the prescription of Brodsky and Lepage [18]. In the approximation when other than two-nucleon degrees of freedom in the deuteron wave function can be neglected, one can unambiguously relate the light-cone deuteron wave functions to those calculated in the rest frame in non-relativistic instant-form calculations [19,20]. The final result for the spectral function in the light-cone approach can be written similarly to Eq.(16), only now S is replaced by the light-cone density matrix: [8,17],

$$S^{LC}(\alpha_s, p_T) = \frac{\sqrt{M^2 + \mathbf{k}^2}}{2 - \alpha_s} |\psi_D(k)|^2, \quad (22)$$

where

$$k \equiv |\mathbf{k}| = \sqrt{\frac{M^2 + p_T^2}{\alpha_s(2 - \alpha_s)} - M^2} \quad (23)$$

is the relative momentum of the two nucleons on the light-cone, and $W_N^{\mu\nu \text{ eff}}$ in Eq.(16) is now the bound nucleon electromagnetic tensor defined on the light-cone. Note that the main operational difference between the instant-form (17) and light-cone (22) impulse approximations is the different relation between the deuteron wave function and the scattering amplitude. Numerical studies have demonstrated that in the kinematical region of interest, $|\mathbf{p}_s| \lesssim 0.5 \text{ GeV}/c$, the difference between the results of the two approximations (for the same deuteron wave function) is quite small — see Sec.III D and Ref. [8].

The structure functions for the scattering from the off-“+”-shell bound nucleons may depend on the variables of this nucleon similarly to the case of the Bethe-Salpeter, or covariant Feynman, approach. In another language this dependence can be interpreted as the presence of non-nucleonic degrees of freedom in the deuteron. With this in mind, we shall use Eq.(16) as the basis for the results discussed in the following sections. Before focusing on specific model calculations of the semi-inclusive deuteron structure functions, however, let us first turn our attention to the validity of the impulse approximation, and the problem of final state interactions in particular.

B. Final State Interactions

In the kinematical region defined by Eq.(9) the contribution of direct processes, where a nucleon is produced in the γ^*N interaction, is negligible [8]. Therefore within the framework of the distorted wave impulse approximation (DWIA) ¹ the total deuteron tensor $W_D^{\mu\nu}$ can be expressed through the nucleon electromagnetic currents as:

$$\begin{aligned} W_D^{\mu\nu}(x, \alpha_s, p_T, Q^2) &\approx \left| \sum \langle D|pn\rangle \langle XN | \mathcal{O}_{IA} + \mathcal{O}_{FSI} | XN \rangle \right|^2 W_N^{\mu\nu}(x, \alpha_s, p_T, Q^2) \\ &\equiv S^{DWIA}(\alpha_s, p_T) W_N^{\mu\nu \text{ eff}}(x, \alpha_s, p_T, Q^2), \end{aligned} \quad (24)$$

where \mathcal{O}_{IA} is the impulse approximation operator, while \mathcal{O}_{FSI} describes the soft final state interactions between the final hadronic products and the spectator nucleon. The function $S^{DWIA}(\alpha_s, p_T)$ now represents the spectral function distorted by FSI effects.

Analysis [21] of the recent high energy deep-inelastic scattering data on slow neutron production [22] is rather indicative that even in heavy nuclei final state interactions are small, so that the average number of hadrons which reinteract in the target does not exceed unity. This indicates that the system which is produced in the γ^*N interaction is quite coherent and interacts at high energies with a relatively small effective cross section,

$$\sigma_{eff} \ll \sigma_{NN}. \quad (25)$$

¹Note that the DWIA approach works best at extreme backward kinematics, Eq.(26) below, where final state interactions have a small contribution, Eq.(27).

Such a situation allows one to factorize the $\gamma^* N$ interaction and use the calculation of FSI in $eD \rightarrow e p n$ processes as a conservative upper limit.

The final simplification which we can gain is to consider the extreme backward kinematics where, in addition to Eq.(9), we also require:

$$\mathbf{p}_T \approx 0. \quad (26)$$

In this case it can be shown [23,24] that:

$$S^{DWIA}(\alpha_s, p_T \approx 0) \sim S(\alpha_s, p_T \approx 0) \left[1 - \frac{\sigma_{eff}(Q^2, x)}{8\pi < r_{pn}^2 >} \frac{|\psi_D(\alpha_s, < p_T >)\psi_D(\alpha_s, 0)|}{S(\alpha_s, p_T \approx 0)/\sqrt{E_s E_s(< p_T^2 >)}} \right], \quad (27)$$

where $< r_{pn}^2 >$ is the average separation of the nucleons within the deuteron, E_s is the spectator nucleon energy, and $E_s(< p_T^2 >) = \sqrt{M^2 + p_z^2 + < p_T^2 >}$ is the energy evaluated at the average transverse momentum $< p_T^2 >^{1/2} \sim 200\text{--}300$ MeV/c transferred for the hadronic soft interactions with effective cross section σ_{eff} . The steep momentum dependence of the deuteron wave function, $|\psi_D(\alpha_s, < p_T >)| \ll |\psi_D(\alpha_s, p_T \approx 0)|$, ensures that FSI effects are suppressed in the extreme backward kinematics, in which case the original impulse approximation expression for $S(\alpha_s, p_T)$ can be used rather than $S^{DWIA}(\alpha_s, p_T)$. Finally, expressing the electromagnetic tensor of the nucleon, $W_N^{\mu\nu}$, through the effective nucleon structure function F_{1N}^{eff} and F_{2N}^{eff} , we can then write for the deuteron tensor:

$$W_D^{\mu\nu} \approx S(\alpha_s, p_T) \left\{ - \left(g_{\mu\nu} + \frac{1}{Q^2} q_\mu q_\nu \right) \frac{1}{M} F_{1N}^{eff} \left(\frac{x}{2 - \alpha_s}, p_T, Q^2 \right) + \left(p_\mu + \frac{p \cdot q}{Q^2} q_\mu \right) \left(p_\nu + \frac{p \cdot q}{Q^2} q_\nu \right) \frac{1}{\nu M^2} F_{2N}^{eff} \left(\frac{x}{2 - \alpha_s}, p_T, Q^2 \right) \right\}. \quad (28)$$

Defining $F_{L,T,TL,TT}^N$ to be the semi-inclusive structure functions in Eqs.(5) with the spectral function factored,

$$F_{L,T,TL,TT}^D = S(\alpha_s, p_T) F_{L,T,TL,TT}^N, \quad (29)$$

one can express the semi-inclusive nucleon functions in terms of the effective structure functions of the nucleon as:

$$F_L^N(x, \alpha_s, p_T, Q^2) = -\sin^2 \delta \frac{\nu}{M} F_{1N}^{eff} \left(\frac{x}{2 - \alpha_s}, p_T, Q^2 \right) + (1 + \cos \delta)^2 \left(\alpha_s + \frac{p \cdot q}{Q^2} \alpha_q \right)^2 \frac{\nu}{\tilde{\nu}} F_{2N}^{eff} \left(\frac{x}{2 - \alpha_s}, p_T, Q^2 \right), \quad (30a)$$

$$F_T^N(x, \alpha_s, p_T, Q^2) = 2F_{1N}^{eff} \left(\frac{x}{2 - \alpha_s}, p_T, Q^2 \right) + \frac{p_T^2}{M^2} \frac{M}{\tilde{\nu}} F_{2N}^{eff} \left(\frac{x}{2 - \alpha_s}, p_T, Q^2 \right), \quad (30b)$$

$$F_{TL}^N(x, \alpha_s, p_T, Q^2) = 2(1 + \cos \delta) \frac{p_T}{M} \left(\alpha_s + \frac{p \cdot q}{Q^2} \alpha_q \right) \frac{\nu}{\tilde{\nu}} F_{2N}^{eff} \left(\frac{x}{2 - \alpha_s}, p_T, Q^2 \right), \quad (30c)$$

$$F_{TT}^N(x, \alpha_s, p_T, Q^2) = \frac{\sin^2 \delta}{2} \frac{p_T^2}{M^2} \frac{\nu}{\tilde{\nu}} F_{2N}^{eff} \left(\frac{x}{2 - \alpha_s}, p_T, Q^2 \right), \quad (30d)$$

where

$$\alpha_q \equiv \frac{\nu - |\mathbf{q}|}{M}, \quad (31a)$$

$$\tilde{\nu} \equiv \frac{p \cdot q}{M} = |\mathbf{q}| \frac{1 + \cos \delta}{2} \alpha_s + \alpha_q \frac{M^2 + p_T^2}{2\alpha_s M}. \quad (31b)$$

Note that within impulse approximation the Callan-Gross relation between the F_{1N}^{eff} and F_{2N}^{eff} structure functions is preserved. Therefore the experimental verification of this relation could serve as another way to identify FSI effects.

Equations (30) and (31) show that at fixed x and $Q^2 \rightarrow \infty$, when $\alpha_q \rightarrow 0$ the longitudinal structure function, F_L^N , does not depend explicitly on the transverse momentum of the nucleon. It contains p_T dependence only in the argument of the bound nucleon structure functions, which arises from the possible nuclear modifications of the nucleon's parton distributions. The above argument, and the fact that FSI of hadronic products with $p_X \gtrsim 1$ GeV/c practically conserve α_s , allows one to conclude that the FSI effect on F_L^N is minimal.

On the other hand the functions F_T , F_{TL} and F_{TT} do explicitly depend on the spectator transverse momentum and therefore the hadronic reinteractions in the final state may strongly affect these structure functions. Such a situation suggests that a separate study of the complete set of the structure functions will allow one to investigate the effects of final state interactions in the deep-inelastic ($e, e'N$) reactions. Note that for the production of spectators with $p_T > 0$, FSI effects are not likely to depend strongly on x for $x > 0.1$. This is because at $x > 0.1$ the essential longitudinal distances in deep-inelastic scattering are small. Therefore for these kinematics one expects that $\sigma_{eff}(Q^2, x) \approx \sigma_{eff}(Q^2)$. The effect of FSI will be investigated in the dedicated deep-inelastic scattering experiments at HERMES which will measure the A -dependence of the forward produced hadrons.

These observations enable us to conclude that in the kinematic region defined by Eqs.(8), (9) and (26), where the contribution of F_L^N is enhanced and FSI effects are small, the differential cross section (4) can be written:

$$\begin{aligned} \frac{d\sigma^{eD \rightarrow epX}}{dx dW^2 d(\log \alpha_s) d^2 \mathbf{p}_T} &\approx \frac{2\alpha_{em}^2}{Q^4} (1-y) \frac{S(\alpha_s, p_T)}{2-\alpha_s} \left[F_L^N + \frac{Q^2}{2\mathbf{q}^2} \frac{\nu}{M} F_T^N \right] \\ &= \frac{2\alpha_{em}^2}{Q^4} (1-y) S(\alpha_s, p_T) F_{2N}^{eff} \left(\frac{x}{2-\alpha_s}, p_T, Q^2 \right) \end{aligned} \quad (32)$$

where we have made the transformation $dQ^2/x \rightarrow dW^2$, with $W^2 = -Q^2 + 2M\nu + M^2$.

Based on the expectation that FSI effects should not strongly depend on x , from Eq.(32) it may be advantageous to consider the ratio of cross sections relative to a given x , in the range 0.1–0.2, where the observed EMC effect in inclusive scattering is small [17,25]:

$$\begin{aligned} G(\alpha_s, p_T, x_1, x_2, Q^2) &\equiv \frac{d\sigma(x_1, \alpha_s, p_T, Q^2)}{dx dW^2 d(\log \alpha_s) d^2 \mathbf{p}_T} \bigg/ \frac{d\sigma(x_2, \alpha_s, p_T, Q^2)}{dx dW^2 d(\log \alpha_s) d^2 \mathbf{p}_T} \\ &= \frac{F_{2N}^{eff}(x_1/(2-\alpha_s), p_T, Q^2)}{F_{2N}^{eff}(x_2/(2-\alpha_s), p_T, Q^2)}. \end{aligned} \quad (33)$$

In our analysis we will consider only production of backward nucleons, $\alpha_s \geq 1$, to suppress contributions from the direct processes where a nucleon is produced in the γ^*N interaction vertex. The more liberal condition, Eq.(9), is in reality sufficient [8].

Note also that for heavier nuclei the FSI becomes much more important in the limit of large x . As was demonstrated in Ref. [17], in this limit rescattering of hadrons produced in the elementary deep-inelastic scattering off the short-range correlation is dynamically enhanced, since the average value of the Bjorken-variable for this mechanism is $\approx x$, as opposed to $x/(2 - \alpha_s)$ for the spectator mechanism. In this sense the deuteron target provides the best way of looking for the EMC effect for bound nucleons.

III. MODELS

In this section we briefly summarize several models of the EMC effect which we use in our analysis and present their predictions for the tagged structure functions. The differences between the models stem from dynamical assumptions about the deformation of the bound nucleon wave functions, and from the fraction of the EMC effect attributed to non-baryonic (mesonic) degrees of freedom in nuclei — from the dominant part in some versions of the binding model, to the models where non-baryonic degrees of freedom play no role, as in the color screening or QCD radiation models. (Other models which have been used in studies of semi-inclusive DIS include the six-quark cluster models discussed in Refs. [26,27].)

A. Binding Models

One of the simplest of the early ideas proposed to explain the nuclear EMC effect was the nuclear binding model, in which the main features of the EMC effect could be understood in terms of conventional nuclear degrees of freedom — nucleons and pions — responsible for the binding in nuclei [17,28–35]. Within the formalism of Sec.II A 1, the inclusive nuclear structure function in the EMC ratio, Eq.(2), is expressed through a convolution of the nuclear spectral function and the structure function of the bound (off-shell) nucleon (c.f. Eq.(16)). Contributions from DIS from the pionic fields themselves, which are needed to balance overall momentum conservation, were considered in Refs. [28,30–32,34], however, their role is most evident only at small x ($x \lesssim 0.2$).

The bulk of the suppression of the EMC ratio (2) at $x \sim 0.6$ in the binding model can be attributed to the fact that the average value for the interacting nucleon light-cone fraction is less than unity. For the case of the deuteron, this corresponds to the average spectator light-cone fraction $\langle \alpha_s \rangle > 1$, which is contrary to what one would have from Fermi motion alone, where the average α_s is < 1 . A relatively minor role is played by the structure function of the bound nucleon itself — the only requirement is that it be a monotonically decreasing function of x [8,28,35]. This is clearly the case for the on-shell structure function, and since the off-shell behavior of the bound nucleon structure function is unknown, most early versions of the binding model simply neglected the possible dependence on p^2 with the expectation that it is not large for weakly bound systems. Only very recently has the issue of off-shell dependence in the bound nucleon structure function been addressed [9,10], where the first attempts to construct models for the p^2 dependence of F_{2N}^{eff} were made. Note that a consequence of assuming the absence of any off-shell effects in F_{2N}^{eff} is that the tagged structure function ratio G in Eq.(33), normalized to the corresponding ratio for a

free proton, would be unity (see Fig.5 below). Any observed deviation of this ratio from unity would therefore be a signal of the presence of nucleon off-shell effects.

If in a dilute system such as the deuteron the nucleon off-shell effects do not play a major role (at least at $x \lesssim 0.7$), one could expand the effective nucleon structure function in a Taylor series about $p^2 = M^2$:

$$F_{2N}^{eff}(x, p^2, Q^2) = F_{2N}(x, Q^2) + (p^2 - M^2) \left. \frac{\partial F_{2N}^{eff}(x, p^2, Q^2)}{\partial p^2} \right|_{p^2=M^2} + \dots \quad (34)$$

Here the off-shell dependence is determined, to order \mathbf{p}^2/M^2 , entirely by the derivative of F_{2N}^{eff} with respect to p^2 . In order to model this correction a microscopic model of the nucleon structure is required [9,10,12]. In any generic quark-parton model, the effective nucleon structure function can be written as an integral over the quark momentum p_q of the quark spectral function ρ :

$$F_{2N}^{eff}(x, p^2, Q^2) = \int dp_q^2 \rho(p_q^2, p^2, x, Q^2). \quad (35)$$

To proceed from Eq.(35) requires additional assumptions about the quark spectral function. The simplest is to assume that the p^2 and p_q^2 dependence in ρ is factored [10,12], which then leads to an explicit constraint on the p^2 derivative of F_{2N}^{eff} from baryon number conservation in the deuteron:

$$\int_0^1 \frac{dx}{x} \left. \frac{\partial F_{2N}^{eff}(x, p^2, Q^2)}{\partial p^2} \right|_{p^2=M^2} = 0. \quad (36)$$

This then allows $\partial F_{2N}^{eff}(x, p^2, Q^2)/\partial p^2$ to be determined from the x -dependence of the on-shell structure function $F_{2N}(x, Q^2)$ in terms of a single parameter, the squared mass of the intermediate state “diquark” system that is spectator to the deep-inelastic collision, $(p-p_q)^2$. One can obtain a good fit to the on-shell nucleon structure function data in terms of this model if one restricts the spectator “diquark” mass to be in the range $(p-p_q)^2 \approx 2-4$ GeV² [10].

A more microscopic model which does not rely on the factorization of the p^2 and p_q^2 dependence in ρ was discussed in Ref. [9]. The quark spectral function there was determined entirely from the dynamics contained in the nucleon–quark–spectator “diquark” vertex function, $\Gamma(p, p_q)$. Within the approximation discussed in Section II A, taking the positive energy nucleon projection only,

$$\rho(p_q^2, p^2, x, Q^2) \rightarrow \text{Tr} \left[(\not{p} + M) \bar{\Gamma}(p, p_q) (\not{p}_q - m_q)^{-1} \not{\epsilon} (\not{p}_q - m_q)^{-1} \Gamma(p, p_q) \right], \quad (37)$$

where m_q is the quark mass. Angular momentum conservation allows two forms for the vertex function Γ , namely scalar and pseudo-vector. In Refs. [9,11,14] it was found that the on-shell data could be well described in terms of only a few of the many possible Dirac structures for Γ . In particular, the vertices were chosen to be $\propto I$ and $\gamma_\alpha \gamma_5$. The momentum dependence of the vertex functions, on the other hand, is more difficult to derive, and must in practice be either parameterized or calculated by solving bound state Faddeev equations in simple models of the nucleon [36]. In order to obtain realistic static properties of the

nucleon, and to account for the bound nature of the nucleon state, the vertex functions must have the form:

$$\Gamma(p, p_q) \propto \frac{(m_q^2 - p_q^2)}{(\Lambda^2 - p_q^2)^n}, \quad (38)$$

where the parameters Λ and n are fixed by comparing with the quark distribution data, and the overall normalization is fixed by the baryon number conservation condition for both the nucleon and deuteron structure functions [9,11,14]. In the comparisons in Sec.III D we use the parameters from the analysis of Ref. [9].

B. Color Screening Model of Suppression of Point-like Configurations in Bound Nucleons

A significant EMC effect in inclusive (e, e') reactions occurs for $x \sim 0.5$ – 0.6 which corresponds to the high-momentum component of the quark distribution in the nucleon. Therefore the EMC effect in this x range is sensitive to a rather rare component of the nucleon wave function where 3 quarks are likely to be close together [17,25]. It is assumed in this model that for large x the dominant contribution to $F_{2N}(x, Q^2)$ is given by the point-like configurations (PLC) of partons which weakly interact with the other nucleons. Note that due to scaling violation $F_{2N}(x, Q^2)$ at $x \gtrsim 0.6$, $Q^2 \gtrsim 10 \text{ GeV}^2$, is determined by the non-perturbative nucleon wave function at $x \gtrsim 0.7$. Thus it is actually assumed that in the nonperturbative nucleon wave function point-like configurations dominate at $x \gtrsim 0.7$. The suppression of this component in a bound nucleon is assumed to be the main source of the EMC effect in inclusive deep-inelastic scattering [17,25]. Note that this suppression does not lead to a noticeable change in the average characteristics of nucleons in nuclei [25].

To calculate the change of the probability of a PLC in a bound nucleon, one can use a perturbation series over a small parameter, κ , which controls corrections to the description of a nucleus as a system of undeformed nucleons. This parameter is taken to be the ratio of the characteristic energies for nucleons and nuclei:

$$\kappa = |\langle U_A \rangle / \Delta E_A| \sim 1/10, \quad (39)$$

where $\langle U_A \rangle$ is the average potential energy per nucleon, $\langle U_A \rangle|_{A \gg 1} \approx -40 \text{ MeV}$, and $\Delta E_A \approx M^* - M \sim 0.5 \text{ GeV}$ is the typical energy for nucleon excitations within the nucleus. Note that $\Delta E_D \gtrsim 2(M_\Delta - M) \sim 0.6 \text{ GeV}$, since the $N - \Delta$ component in the deuteron wave function is forbidden due to the zero isospin of the deuteron.

To estimate the deformation of the bound nucleon wave function we consider a model where the interaction between nucleons is described by a Schrödinger equation with potential $V(R_{ij}, y_i, y_j)$ which depends both on the inter-nucleon distances (spin and isospin of nucleons) and the inner variables y_i and y_j , where y_i characterizes the quark-gluon configuration in the i -th nucleon [17,25,37]. The Schrödinger equation can be represented as:

$$\left[-\frac{1}{2m_N} \sum_i \nabla_i^2 + \sum_{i,j} V(R_{ij}, y_i, y_j) + \sum_i H_0(y_i) \right] \psi(y_i, R_{ij}) = E \psi(y_i, R_{ij}). \quad (40)$$

Here $H_0(y_i)$ is the Hamiltonian of a free nucleon. In the nonrelativistic theory of the nucleus the inter-nucleon interaction is averaged over all y_i . Thus the nonrelativistic $U(R_{ij})$ is related to V as:

$$U(R_{ij}) = \sum_{y_i, y_j, \tilde{y}_i, \tilde{y}_j} \langle \phi_N(y_i) \phi_N(y_j) | V(R_{ij}, y_i, y_j, \tilde{y}_i, \tilde{y}_j) | \phi_N(\tilde{y}_i) \phi_N(\tilde{y}_j) \rangle, \quad (41)$$

where $\phi_N(y_i)$ is the free nucleon wave function. The unperturbed wave function is the solution of Eq.(40) with the potential V replaced by U . Treating $(U - V)/(E_i - E_N)$ as a small parameter, where E_i is the energy of an intermediate excited nucleon state, one can calculate the dependence of the probability to find a nucleon in a PLC to the momentum of a nucleon inside the nucleus. One finds that this probability is suppressed as compared to the similar probability for a free nucleon by the factor [25]:

$$\delta_A(\mathbf{k}^2) \approx 1 - 4(\mathbf{k}^2/2M + \epsilon_a)/\Delta E_A, \quad (42)$$

where $\Delta E_A = \langle E_i - E_N \rangle \approx M^* - M$, in first order of the perturbation series. An estimate of higher order terms gives [17]:

$$\delta_A(\mathbf{k}^2) = (1 + z)^{-2}, \quad z = (\mathbf{k}^2/M + 2\epsilon_a)/\Delta E_A. \quad (43)$$

The x dependence of the suppression effect is based on the assumption that the PLC contribution in the nucleon wave function is negligible at $x \lesssim 0.3$, and gives the dominant contribution at $x \gtrsim 0.5$ [25,38]. We use a simple linear fit to describe the x dependence between these two values of x [38]. One can then obtain an estimate for R_A in Eq.(2) for large A at $x \sim 0.5$,

$$R_A(x) |_{x \sim 0.5} \sim \delta_A(\mathbf{k}^2) \approx 1 + \frac{4\langle U_A \rangle}{\Delta E_A} \sim 0.7 - 0.8, \quad (44)$$

since here Fermi motion effects are small. The excitation energy ΔE_A for the compressed configuration is estimated as $\Delta E_A \sim (M(1400) - M) - (M(1680) - M) \sim (0.5-0.8)$ GeV, while $\langle U_A \rangle \approx -40$ MeV. Since $\langle U_A \rangle \sim \langle \rho_A(r) \rangle$ the model predicts also the A dependence of the EMC effect, which is consistent with the data [17].

For the deuteron target we can deduce from Eq.(44) using $\langle U_D \rangle / \langle U_{Fe} \rangle \sim 1/5$,

$$R_D(x, Q^2) |_{x \sim 0.5} \approx 0.94 - 0.96. \quad (45)$$

This number may be somewhat overestimated because, as discussed above, due to the isoscalarity of D , low-energy excitations in the two-nucleon system are forbidden, leading to $\Delta E_D > \Delta E_A$.

This model represents one of the extreme possibilities that the EMC effect is solely the result of deformation of the wave function of bound nucleons, without attributing any extra momentum to be carried by mesons. A distinctive feature of this explanation of the EMC effect is that the deformation of the nucleon should vary with inter-nucleon distances in nuclei (with nucleon momentum in the nucleus). For $|\mathbf{k}| \sim 0.3 - 0.4$ GeV/c the deviation from the conventional quantum-mechanical model of a deuteron is expected to be quite large (factor ~ 2). Actually, the size of the effects may depend not on \mathbf{k}^2 only but on p_T and α_s separately, because the deformation of a bound nucleon may be more complicated than suggested by this simple model.

C. QCD Radiation, Quark Delocalization.

It was observed in Refs. [39,40] that the original EMC data could be roughly fitted as:

$$\frac{1}{A}F_{2A}(x, Q^2) = \frac{1}{2}F_{2D}\left(x, Q^2\xi_A(Q^2)\right), \quad (46)$$

with $\xi_{Fe}(Q^2) \approx 2$ for $Q^2 \approx 20 \text{ GeV}^2$, the so-called dynamical rescaling. The phenomenological observation has been interpreted as an indication that gluon radiation occurs more efficiently in a nucleus than in a free nucleon (at the same Q^2) due to quark delocalization, either in a bound nucleon (or in two nearby nucleons) [39,41–43] or in the nucleus as a whole [41].

The Q^2 dependence of $\xi(Q^2)$ follows from the requirement that both sides of Eq.(46) should satisfy the QCD evolution equations. In the leading logarithmic approximation one has:

$$\xi_A(Q^2) = \xi_A(Q_0^2)^{\alpha_{QCD}(Q_0^2)/\alpha_{QCD}(Q^2)}. \quad (47)$$

Experimentally $d \ln F_{2D}(x, Q^2)/d \ln Q^2$ is positive if $x > x_0$ and negative if $x < x_0$, where $x_0 = 0.15 \pm 0.05$. So Eq.(46) predicts that the EMC effect should vanish at $x \approx x_0$.

If the confinement size in a nucleus, λ_A , is larger than that in a free nucleon, one may expect that the Q^2 evolution of the parton distributions in nuclei (the bremsstrahlung of gluons and quarks) may start at $Q_0^2(A) < Q_0^2(N)$. To reproduce Eq.(46) one should have:

$$Q_0^2(A)/Q_0^2(N) = [\xi(Q_0^2(A))]^{-1}. \quad (48)$$

Assuming on dimensional grounds that the radii of quark localization in a nucleus, λ_A , and in a nucleon, λ_N , are related via:

$$Q_0^2(A)\lambda_A^2 = Q_0^2(N)\lambda_N^2 \quad (49)$$

to the scale for the onset of evolution, Q_0^2 , one finally obtains:

$$\lambda_A/\lambda_N \approx \xi(Q_0^2(A))^{1/2}. \quad (50)$$

A fit to the original EMC data using Eq.(46) leads to $\xi_{Fe}(20 \text{ GeV}^2)=2$. For $\mu_{Fe}^2 = 0.67 \text{ GeV}^2$ and $\Lambda_{\overline{MS}} \approx 250 \text{ MeV}$ this corresponds to:

$$\lambda_{Fe}/\lambda_D \approx 1.15 \quad \lambda_{A \sim 200}/\lambda_D \approx 1.19. \quad (51)$$

Since in this model the delocalization is approximately proportional to the nuclear density, one can expect that the effect is also proportional to the \mathbf{k}^2 of a bound nucleon. Fixing the parameters of the model to fit the Fe data, and assuming $\lambda(k)/\lambda = 1 + a\mathbf{k}^2$ we obtain $a \approx 0.4/ < \mathbf{k}^2 >_{Fe}$, where $< \mathbf{k}^2 >_{Fe} \sim 0.08 \text{ GeV}^2/c^2$. Using this expression for $\lambda(k)$ one can calculate the $\xi_d(Q^2, k)$ according to Eqs.(47)–(50) and express the effective structure function of a bound nucleon as:

$$F_{2N}^{eff}(x, \alpha_s, p_T, Q^2) = F_{2N}\left(x, Q^2\xi(Q^2, k)\right), \quad (52)$$

where k is defined through α_s and p_T according to Eq.(23).

D. Numerical estimates

Comparison of predictions of the models for the nuclear EMC effect considered in the preceding Sections is most meaningful in the kinematic range where, firstly, FSI effects are small, and secondly, the instant-form and light-cone prescriptions for the deuteron spectral function lead to similar results.

Direct calculation of the FSI contribution to the cross section would require knowledge of the full dynamics of the final N - X system, which is a practically impossible task given the present level of understanding of nonperturbative QCD. However, it is possible to estimate the uncertainty which would be introduced through neglect of FSI, by using the calculation of FSI effects in the high-energy $d(e, e'p)n$ (break-up) reaction in Refs. [23,24], and replacing the p - n rescattering cross section by an effective cross section for the $p(n)$ - X interaction, Eq.(27). For the effective cross section σ_{eff} one can use the results of the recent analysis [21] of soft neutron production in the high-energy deep-inelastic scattering of muons from heavy nuclei [22], which yielded an upper limit of $\sigma_{eff} \approx 20$ mb. To be on the conservative side, in the following estimates we therefore use the value of $\sigma_{eff} \approx 20$ mb suggested by string models of FSI (for a recent summary see Ref. [44]). Furthermore, by retaining only the imaginary part of the spectator nucleon rescattering amplitude, one obtains an upper limit of the FSI effect, since the real part will contribute to elastic rescattering only, effectively suppressing the value of σ_{eff} . In Fig.1 we illustrate the α_s dependence of the ratio of the (light-cone) spectral function including FSI effects within the DWIA, Eq.(27), to that calculated without FSI effects. At extreme backward kinematics ($p_T \approx 0$) one sees that FSI effects contribute less than $\sim 5\%$ to the overall uncertainty of the $d(e, e'N)X$ cross section for $\alpha_s \lesssim 1.5$. As mentioned above, this number can be considered rather as an upper limit on the uncertainties due to FSI. At larger p_T ($\gtrsim 0.3$ GeV/c), and small α_s (≈ 1), the double scattering contribution (which is not present for the extreme backward case in Eq.(27)) plays a more important role in FSI [23]. Because its sign is positive, it tends to cancel some of the absorption effects of FSI at large p_T (for a detailed discussion of the double scattering contribution in FSI see Ref. [23]). Note also that the FSI effects do not change significantly with energy for fixed Q^2 . Thus, for the ratios discussed, where the cross sections are compared for the same Q^2 but different x , the changes due to FSI effects are even smaller.

To extract unambiguous information from the semi-inclusive cross section ratios on the medium modifications of the nucleon structure discussed in this Section requires one to establish the regions of kinematics where the differences between the various prescriptions for the deuteron spectral function are minimal. In Fig.2 we illustrate the α_s and p_T dependence of the ratio of spectral functions calculated in the light-cone (22) and instant-form (17) approaches. For $p_T \leq 0.1$ GeV/c the light-cone and instant-form predictions differ up to the 20% for the entire range of $\alpha_s \leq 1.5$. However, choosing the isolated values of $\alpha_s \leq 1.2$ or $\alpha_s \approx 1.4$ one can confine the uncertainty in the spectral function to within 10%, which will offer the optimal conditions in which to study the nucleon structure modification.

In Fig.3 we show the α_s dependence of the ratio of the effective proton structure function, F_{2p}^{eff} , for extreme backward kinematics, $p_T = 0$, to the structure function of a free proton. As expected, the suppression of the ratio in the version of the binding model with explicit nucleon off-shell corrections is quite small, $\lesssim 10\%$ for $\alpha_s < 1.5$, reflecting the relatively minor

role played here by off-shell effects in the nucleon structure function. (Note that in versions of the binding model in which there is no p^2 dependence in the effective nucleon structure function this ratio would be unity.) The effects in the PLC suppression and rescaling models in Fig.3 are somewhat larger. For a neutron target one predicts similar results, however, the neutron structure function is not as well known experimentally as the proton due to the absence of free neutron targets (see Section IV). For this reason we restrict our discussion to ratios of proton structure functions only.

In Fig.4 we illustrate the dependence of F_{2p}^{eff}/F_{2p} on the variable $(p^2 - 2M\epsilon)/M^2$, where p^2 is the bound nucleon virtuality defined in Eq.(20) and $\epsilon = -2.2$ MeV is the deuteron binding energy. The comparison in Fig.4 is done for different values of x and α_s . It is noticeable that because of the negligible amount of the PLC component in the nucleon wave function at $x \lesssim 0.3$, the PLC suppression model predicts no modification of the structure function in this region. On the other hand, at $x = 0.6$ it predicts a maximal effect because of the PLC dominance in the nucleon wave function here.

Because the momentum (virtuality) dependent density effect generates the modification of the bound nucleon structure functions in the PLC suppression and rescaling models, the ratios for these models at $x \gtrsim 0.5$ – 0.6 vary similarly with α_s and $(p^2 - 2M\epsilon)/M^2$, as does also the off-shell model ratio. However the mechanism for such a variation is different, which is clearly seen in Fig.5, where the x dependence of the same ratio is represented at different values of α and fixed $p_T = 0$. The curves for the off-shell model extend only to $x \sim 0.7$ because for larger x values the approximations discussed in Sec.II A 1 involved in obtaining Eq.(16) become numerically less justified [9,11].

To further reduce any uncertainties due to the deuteron spectral function in the model comparisons, we concentrate on the predictions for the ratio $G(\alpha_s, p_T, x_1, x_2, Q^2)$, defined in Eq.(33), of experimentally measured cross sections at two different values of x [25]. Since the function G is defined by the ratio of cross sections at the same α_s and p_T , any uncertainties in the spectral function cancel. This allows one to extend this ratio to larger values of α_s , thereby increasing the utility of the semi-inclusive reactions when analyzed in terms of this function. Figure 6 shows the α_s and Q^2 dependence of $G(\alpha_s, p_T, x_1, x_2, Q^2)$ at $p_T = 0$. The values of x_1 and x_2 are selected to fulfill the condition $x_1/(2 - \alpha_s) = 0.6$ (large EMC effect in inclusive measurements) and $x_2/(2 - \alpha_s) = 0.2$ (essentially no EMC effect in inclusive measurements). Again, the PLC suppression and Q^2 rescaling models predict a much faster drop with α_s than does the binding/off-shell model, where the α_s dependence is quite weak.

IV. EXTRACTION OF THE NEUTRON/PROTON STRUCTURE FUNCTION RATIO

The presence of an EMC effect in the deuteron leads to substantial suppression of the deuteron structure function at large x compared to what one would expect from models without non-nucleonic degrees of freedom. For example, the estimate of Ref. [25],

$$\frac{F_{2D}(x, Q^2)}{F_{2p}(x, Q^2) + F_{2n}(x, Q^2)} \approx \frac{1}{4} \frac{2F_{2A}(x, Q^2)}{AF_{2D}(x, Q^2)} \Big|_{A \sim 60; 0.3 < x < 0.7} \quad (53)$$

which is valid for a rather wide class of models in which the EMC effect is proportional to the mean value of p^2 in nuclei, gives a ratio ~ 3 – 5% below unity at $x \sim 0.6$ – 0.7 . Calculations

of the deuteron structure function in models in which binding effects are explicitly taken into account also produce similar effects [11].

Inclusion of the EMC effect in the extraction of the the neutron structure function from the inclusive eD scattering data leads to significantly larger values for F_{2n}/F_{2p} than the $1/4$ value obtained in analyses in which only Fermi motion is included. The values for F_{2n}/F_{2p} with inclusion of the EMC effect at $x \sim 0.6$ [45–47] are in fact much closer to the expectation of $3/7$ from perturbative QCD, predicted by Farrar and Jackson [48]. Therefore observation of a value of F_{2n}/F_{2p} higher than the $1/4$ extracted from inclusive data in the early analyses would by itself serve as another proof of the presence of the EMC effect in the deuteron.

Although other methods to obtain the large- x n/p ratio (or the d/u ratio) have been suggested, none has so far been able to clearly discriminate between the different $x \rightarrow 1$ limits for F_{2n}/F_{2p} (namely, $d/u \rightarrow 0$, which is the minimal possible value allowed in the parton model, which corresponds to $F_{2n}/F_{2p} \rightarrow 1/4$, and $d/u \rightarrow 1/5$ in perturbative QCD [48]). For example, with ν and $\bar{\nu}$ beams on proton targets one can in principle measure the u and d quark distributions separately, however, the statistics in ν experiments in general are relatively poor. Another possibility would be to extract the d/u ratio from charged lepton asymmetries at large rapidities, in W -boson production in $p\bar{p}$ scattering [49], although here it may be some time still before a sufficient quantity of large-rapidity events at the CDF at Fermilab are accumulated.

On the other hand, with tagged deuteron experiments planned for HERMES, a study of the tagged structure functions may allow a resolution of this ambiguity [6,8,50]. It is important that the ratio of tagged structure functions interpolated to the nucleon pole should be exactly equal to the free nucleon ratio — this is the analog of the Chew-Low interpolation for the pion case:

$$\frac{F_{2n}(x, Q^2)}{F_{2p}(x, Q^2)} \approx \frac{F_{2n}^{eff}(x/(2 - \alpha_s), p_T, Q^2)}{F_{2p}^{eff}(x/(2 - \alpha_s), p_T, Q^2)} \Big|_{\alpha_s \approx 1, p_T \approx 0}. \quad (54)$$

In practice the data cannot be accumulated for too small \mathbf{p} . However, we observed above that deviations of the ratio from the free limit is proportional to p^2 with a good accuracy. Hence, if one samples the data as a function of p^2 interpolation to the pole $p^2 - M^2 = 2M\epsilon$ should be smooth. In practice, considering momentum intervals of 100–200 MeV/c and 200–350 MeV/c would be sufficient. A potential problem with Eq.(54) is that at very large x ($x \gtrsim 0.7$) the factorization approximation itself breaks down and higher order corrections to Eq.(16), which are $\propto \mathbf{p}^4$, must be included if one wants accuracy to within a few %. To be on the safe side one should therefore restrict the analysis to smaller spectator momenta, below 200 MeV/c.

V. CONCLUSION

Despite the many years of study of the deviations from unity of the ratios of nuclear to deuteron cross sections in inclusive high-energy scattering, we have been unable to isolate the microscopic origin of the nucleon structure modification in the nuclear environment [1–5] — the effect can be described in terms of a number of models based on quite disparate physical assumptions. In this paper we have argued that by allowing greater accessibility to kinematic

variables not available in inclusive reactions, semi-inclusive deep-inelastic processes offer the possibility to make further progress in understanding the origin of the nuclear EMC effect.

In particular, measurements of tagged structure functions of the deuteron can probe the extent of the deformation of the intrinsic structure of a bound nucleon. Taking ratios of semi-inclusive cross sections at different values of x enables the cancellation in the tagged structure function ratio, Eq.(33), of the dependence on the deuteron wave function, thus permitting the nucleon structure to be probed directly. Our results show that possible contamination of the signal due to the final state interactions of the spectator nucleon with the hadronic debris can be minimized by tagging only on the slow backward nucleons in the target fragmentation region. This may then allow one to discriminate between models in which the EMC effect is attributed entirely to the nucleon wave function deformation, and ones in which the effect arises from more traditional descriptions in terms of meson–nucleon degrees of freedom associated with nuclear binding.

As a by-product of the semi-inclusive measurements, one may also be able to extract information on the large- x behavior of the neutron to proton structure function ratio, by detecting recoiling protons and neutrons with small transverse momentum in the extreme backward kinematics. Extracting this ratio from inclusive ep and eD data is fraught with large uncertainties arising from different treatments of the nuclear physics in the deuteron. Observation of an asymptotic value for F_{2n}/F_{2p} which is larger than the ‘canonical’ $1/4$ would in itself be proof of the presence of an EMC effect in the deuteron.

First information about the spectra of backward nucleons in eD scattering is likely to come from the Jefferson Lab experiment # 94-102 [6], and from the HERMES experiment at HERA, where the necessary counting rate may be achieved after an upgrade of the detector [51]. Having two energy ranges would be very useful for checking the basic production mechanism, and understanding backgrounds and corrections due to final state interactions, which are likely to depend substantially on the incident energy. Final state interaction effects can also be tracked by studying the production of nucleons from heavier nuclei. Probing the quark-gluon structure of short-range correlations with heavy nuclei targets should further enable one to determine whether nucleon deformations predominantly depend on the nucleon momentum, or also on A .

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FIGURES

FIG. 1. The α_s dependence of the ratio of cross sections calculated with FSI effects within the DWIA, and without FSI effects. The curves correspond to different values of the spectator nucleon transverse momentum (in GeV/c).

FIG. 2. The α_s dependence of the ratio of cross sections calculated within the light-cone and instant-form approaches. The curves correspond to different values of the spectator nucleon transverse momentum (in GeV/c).

FIG. 3. The α_s dependence of F_{2p}^{eff}/F_{2p} for $x = 0.6$ and $p_T = 0$. Dashed line is the PLC suppression model, dotted is the rescaling model, and dot-dashed the binding/off-shell model.

FIG. 4. The $(p^2 - 2M\epsilon)/M^2$ dependence of F_{2p}^{eff}/F_{2p} , for $\alpha_s = 1.2$ and 1.4 , and $x = 0.3$ and 0.6 . Curves are as in Fig.3.

FIG. 5. The x dependence of F_{2p}^{eff}/F_{2p} for $\alpha_s = 1.2$ and 1.4 , with $p_T = 0$. Curves are as in Fig.3.

FIG. 6. The α_s dependence of $G(\alpha_s, p_T, x_1, x_2, Q^2)$, with $x_1 = x/(2 - \alpha_s) = 0.6$ and $x_2 = x/(2 - \alpha_s) = 0.2$, for $p_T = 0$. $G^{eff}(\alpha_s, p_T, x_1, x_2, Q^2)$ is normalized to $G^{eff}(\alpha_s, p_T, x_1, x_2, Q^2)$ calculated with the free nucleon structure function. Curves are as in Fig.3.

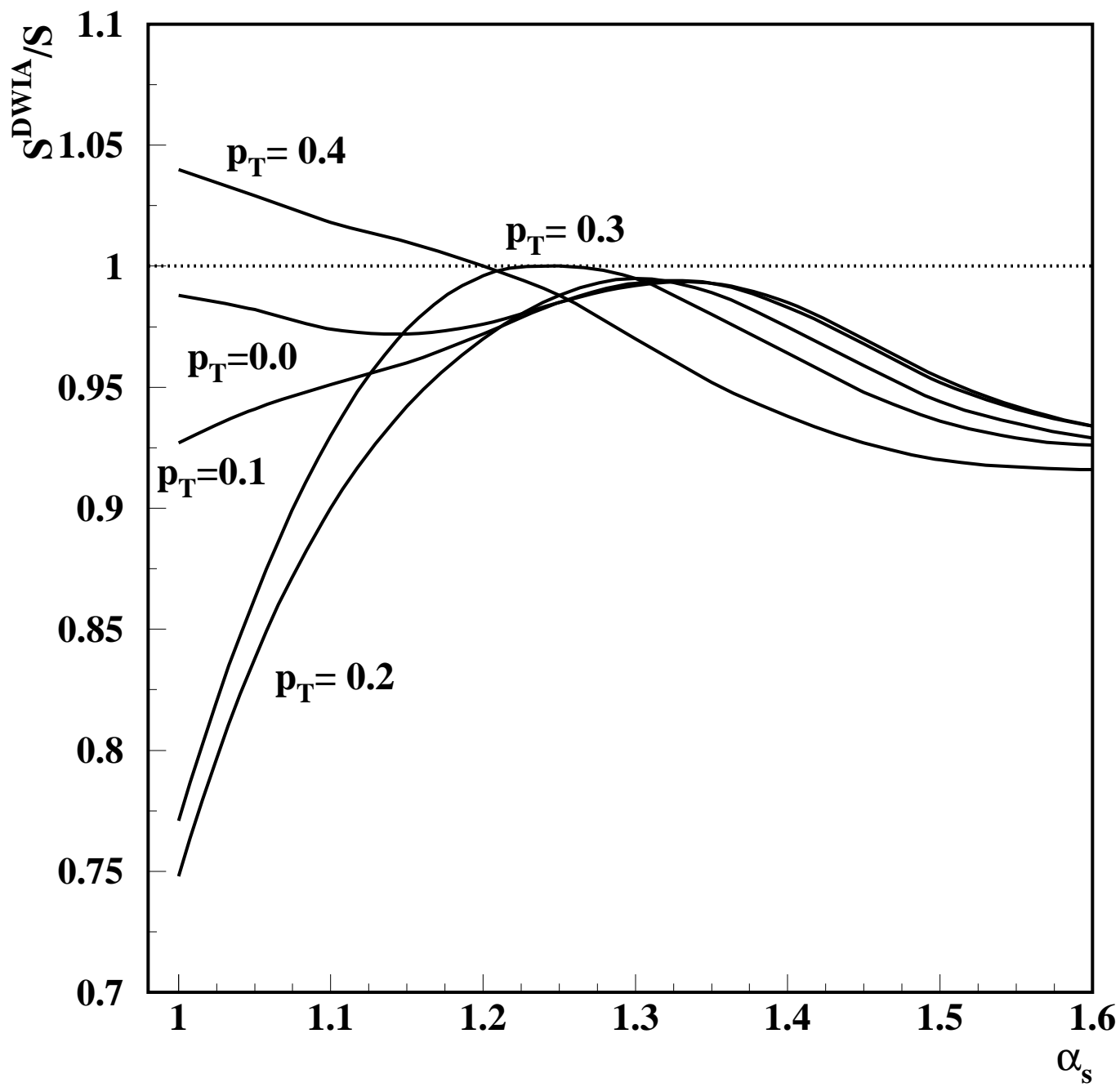


Fig.1

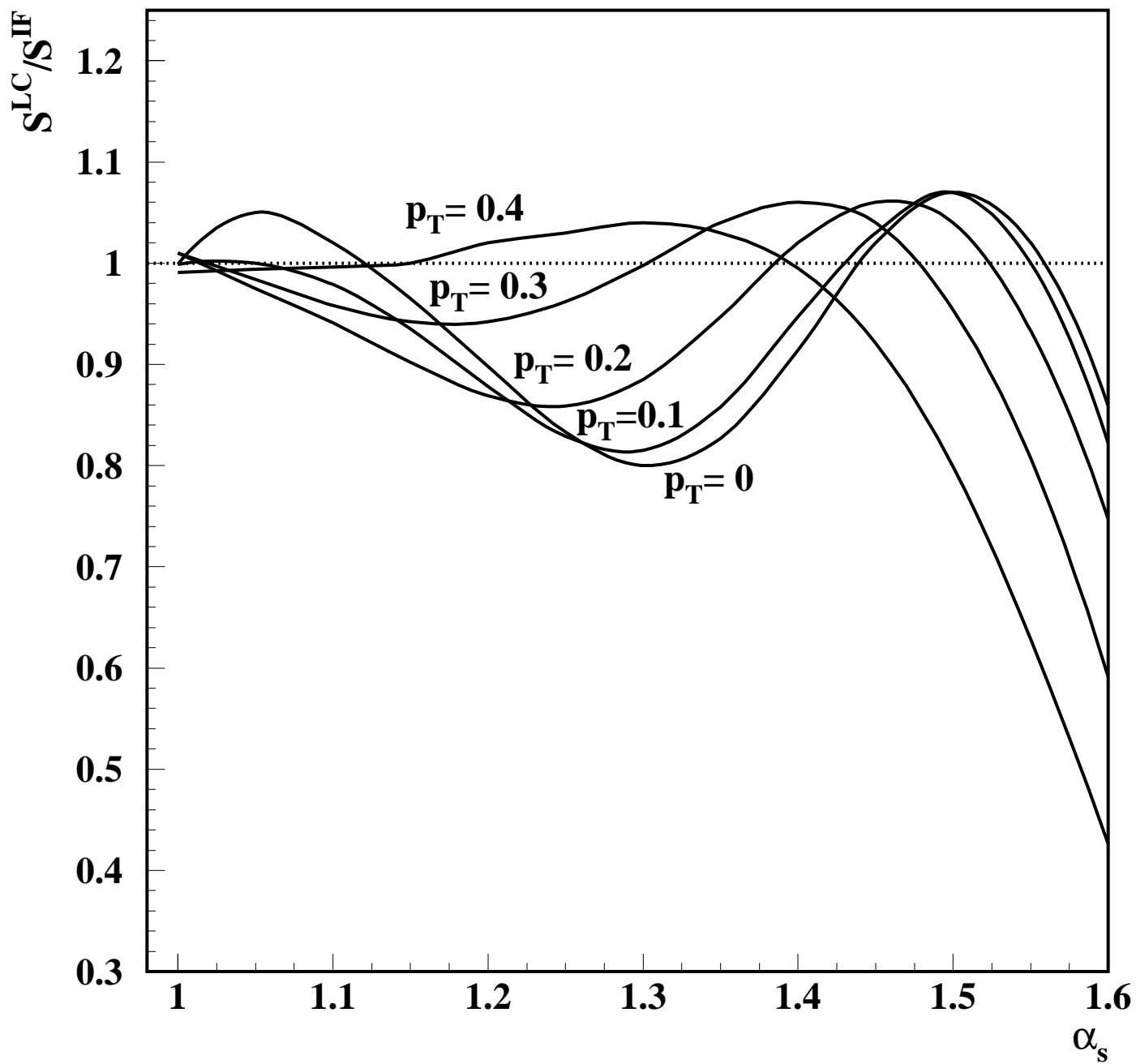


Fig.2

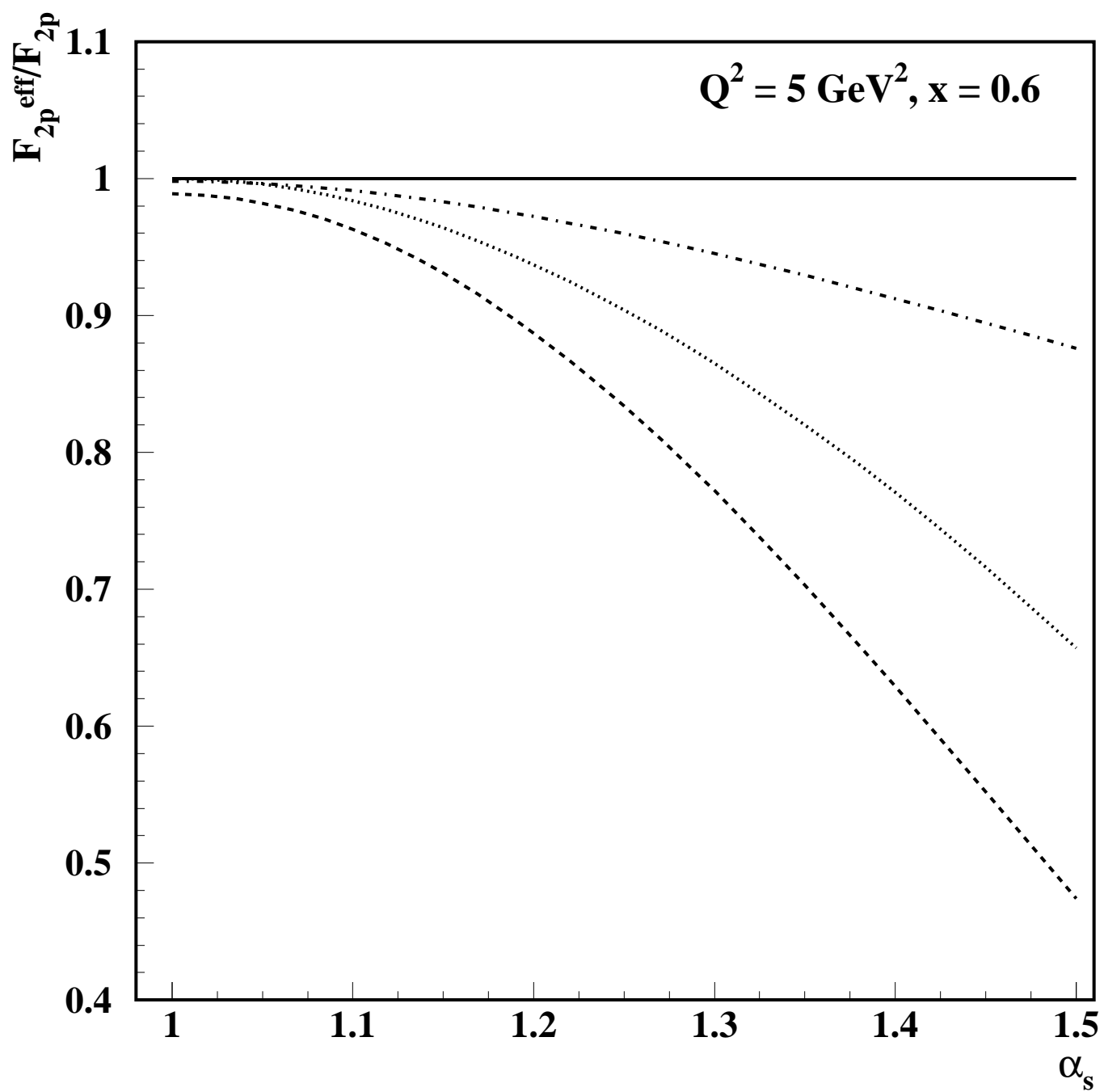


Fig.3

$$Q^2 = 5 \text{ GeV}^2$$

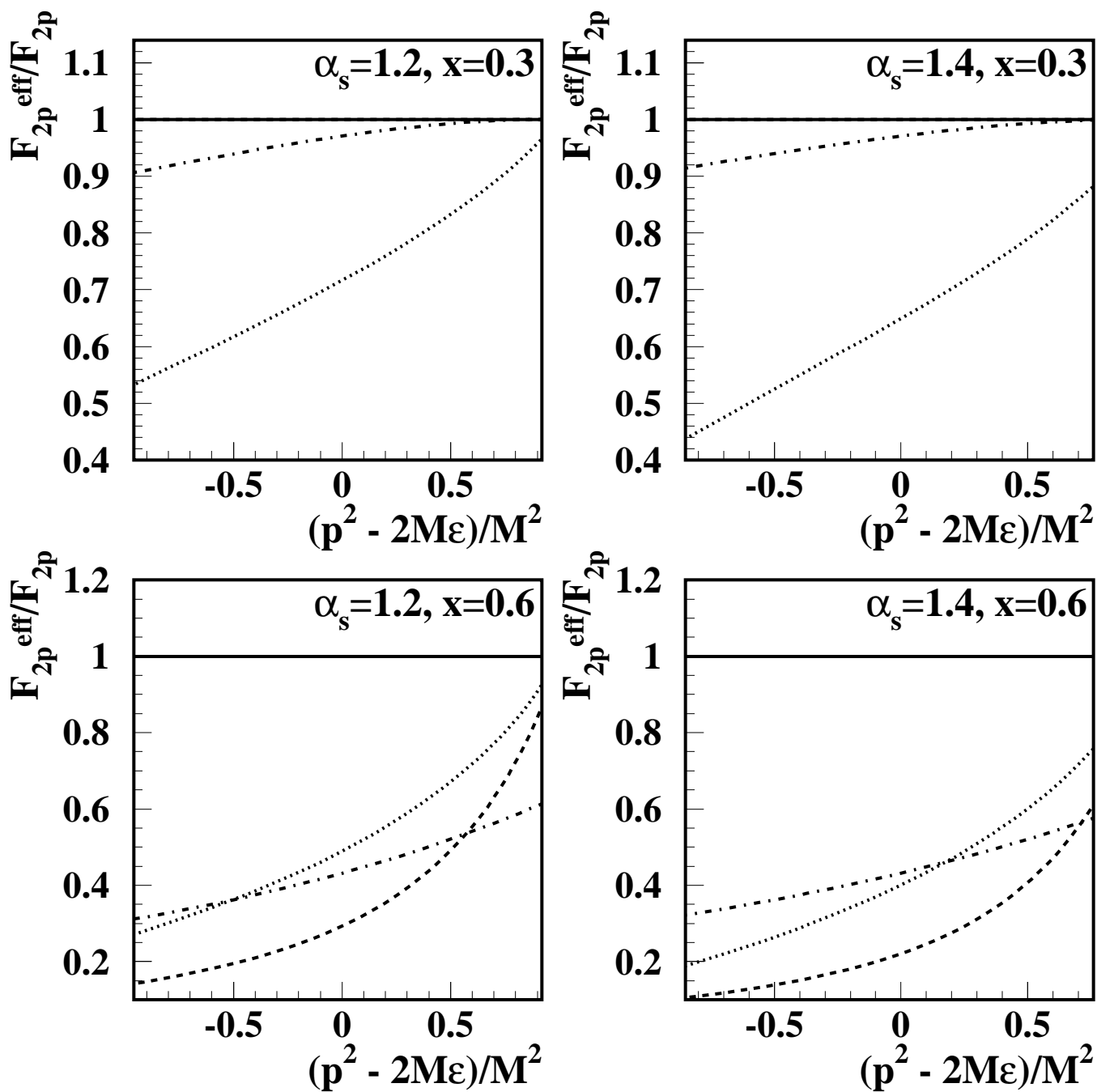


Fig.4

$$Q^2 = 5 \text{ GeV}^2$$

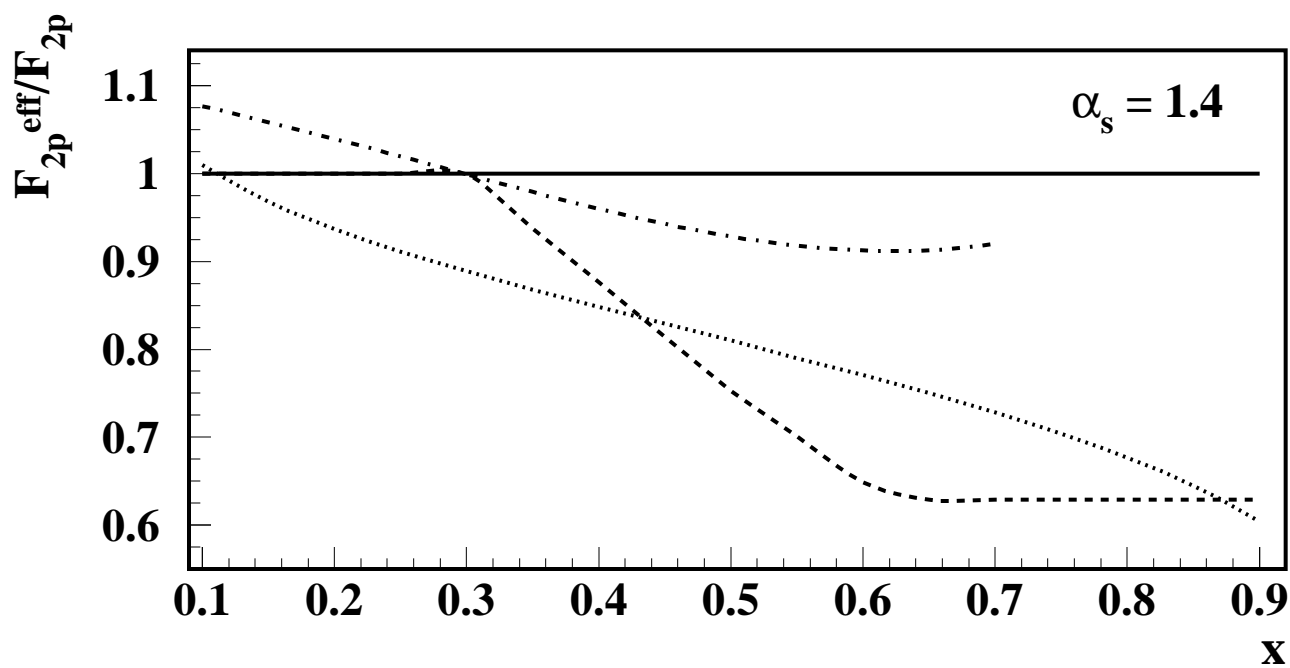
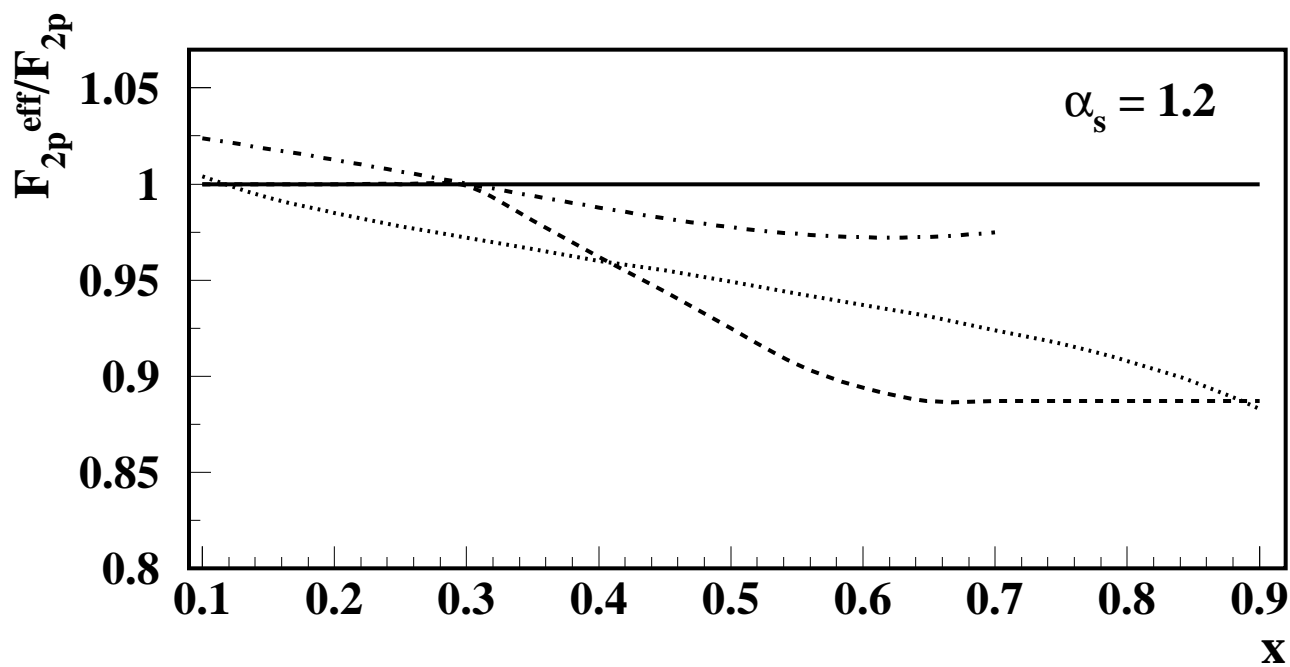


Fig.5

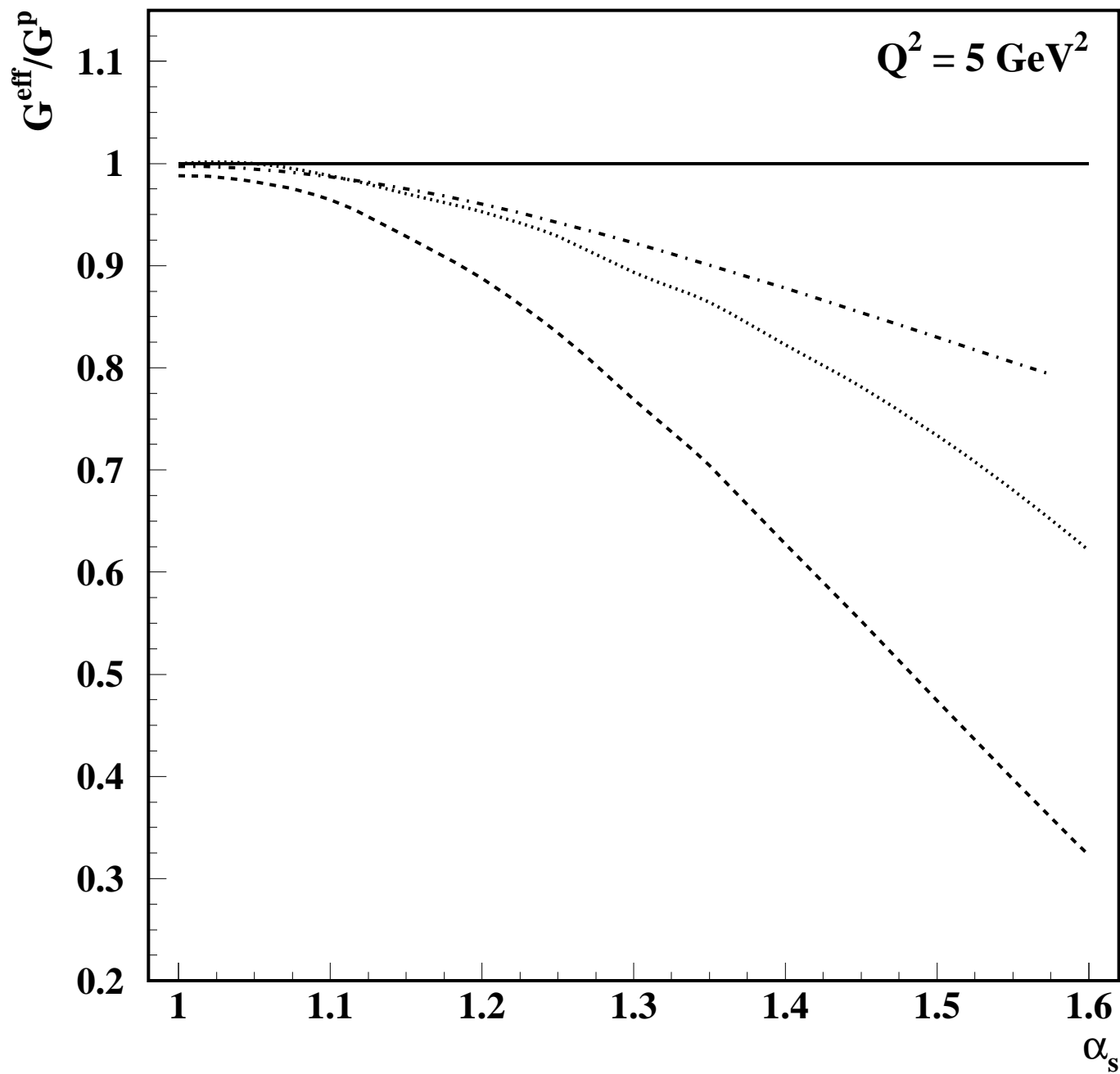


Fig.6